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Ballistic transport of composite fermions in narrow cross junctions

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Abstract. The transport properties of the half-filled Landau level in crossed-wire junctions are presented. A dip in the bend resistance due to ballistic transmission of composite fermions has been observed near $\nu = 1/2$ and 1/4. The temperature dependence of the dips is found to be strong due to the small Fermi energy of the composite fermions. The dips are found to shift to lower magnetic fields, indicating that the effective magnetic field for the composite fermions in the narrow channel is not uniform.

1. Introduction

The fractional quantum Hall effect (FQHE) [1] is observed in a high mobility two-dimensional electron gas (2DEG) at low temperatures and very high magnetic fields. The longitudinal resistance shows a rich spectrum of minima and the Hall resistance R_H is quantized at $(h/e^2)/\nu$, where the Landau level filling factor $\nu = p/q$ with p and q being integral and q being odd. Despite the similar appearance of the integer quantum Hall effect (IQHE) and the FQHE, strong Coulomb interactions between electrons are responsible for the FQHE [2, 3].

Recently, Jain [4] has provided a new interpretation of the FQHE based on the idea of the 'composite fermion' (CF), for which an even number $2s$ of magnetic flux quanta are attached to each electron. Since the CF experiences the external magnetic field B and the magnetic flux attached to the other CFs, it sees, in the mean-field approximation, an effective magnetic field:

$$\Delta B = B \pm 2sn_s\phi_0 \tag{1}$$

where $\phi_0 = h/e$ is the magnetic flux quantum and n_s is the density of the CFs which equals that of the electrons. In this approach, the FQHE appears to be simply the IQHE of the CFs. The fictitious magnetic field exactly cancels the external field at $\nu = 1/2s$. Halperin, Lee and Read [5] have shown that the system of interacting electrons in high magnetic fields can be replaced by a system of independent CFs in zero effective magnetic field. Consequently, the system is anticipated to behave as a metallic state with a well defined Fermi level [5]. The resistance oscillation in the fractional quantum Hall regime can be regarded as the Shubnikov–de Haas oscillation. The effective mass and the single-particle scattering time of the CFs were determined by conventional magnetotransport analysis [6]. Du *et al* [7] reported that the CF mass diverges as $\nu \to 1/2$, indicating that the half-filled state is not an ordinary Fermi liquid. Moreover, possible experimental consequences arising from the existence of the CFs have become apparent. The expected behaviours of the CF have been demonstrated in recent experiments [8–10].

The traditional theory [2, 3] does not predict a quantum Hall state in a 2DEG at $\nu = 1/2$, and it has never been observed experimentally except in certain double-layer systems [11]. However, Timp *et al* [12] reported the observation of a plateau in R_H near $\nu = 1/2$ in narrow crossed-wire junctions. Although they ascribed it to a fractional quantum Hall state, which may be realized in narrow wires, the mechanism remained unclear since the available theory [13] did not include the invasive coupling of the voltage leads. The theory [13] has not been confirmed experimentally so far. The CF picture of the half-filled Landau level provides a unique possibility of explaining the plateau at $\nu = 1/2$ observed in narrow cross junctions in terms of the ballistic transport of the CFs. Typical ballistic transport phenomena of classical electrons in narrow crossed-wire junctions are the quenching of R_H [14] and the negative bend resistance R_B [15] at low magnetic fields.

In this paper, we investigate the transport properties of the half-filled Landau level in narrow cross junctions in order to examine the validity of the CF picture in microstructures. Previous experimental investigations [9, 10] of CFs have concentrated on the ballistic transport properties in the two-dimensional regions of the samples. The effects of the gauge fluctuations, which may partly destroy the Fermi liquid description [5], may be enhanced in quasi-one-dimensional systems as effects of the quantum fluctuation generally are of significant importance in lower-dimensional systems. We find that pronounced dips appear in the bend resistance near $\nu = 1/2$, 1/4, and 3/4 as the width of the wire is reduced. The composite fermion mass has been reported to diverge very strongly as $\nu \to 1/2$. The previous experiments [6–10] examined the CFs in the vicinity of $\nu = 1/2$, whereas the negative bend resistance occurs exactly at $\nu = 1/2$. If the composite fermion effective mass is sensitive even to the presence of small effective magnetic fields, the bend resistance may have the potential to detect it. Therefore, the negative bend resistance is expected to provide information about the non-ordinary Fermi liquid at $\nu = 1/2$. It is shown that the dip exhibits temperature dependence in our measurement temperature range of down to 0.3 K despite the fact that the negative bend resistance due to electrons is robust against temperature and that the CF state in a 2DEG becomes temperature independent at low temperatures.

2. Samples and experimental details

The samples were made from a high-mobility GaAs/$Al_xGa_{1-x}As$ heterostructure with the spacer layer thickness of 80 nm grown by molecular beam epitaxy. The heterointerface is located 150 nm below the surface. The mobility and the electron density in the starting material at $T = 0.3$ K are $\mu = 57$ m^2 V^{-1} s^{-1} and $n_s = 0.81 \times 10^{15}$ m^{-2}, respectively, in the dark. After illuminations with a light emitting diode, the mobility is improved to $\mu = 10^5$ m^2 V^{-1} s^{-1} at $n_s = 1.8 \times 10^{15}$ m^{-2}. The corresponding transport mean free path l_0^{el} of electrons is estimated to be 2.5–7 μm. Although the mobility of our heterostructure is not as large as that in [9] and [10], the length of required ballistic trajectories is much smaller for the negative bend resistance. The ballistic transport effect is hence observable in our samples as will be shown below. The crossed-wire junctions with several widths were fabricated using electron beam lithography. As shown in figure 1, the device consists of four narrow constrictions with several different lithographic constriction widths W_{lith} between 0.8 and 4.0 μm. The effective conduction width is narrower than W_{lith} due to the side-wall depletion. To avoid process-induced damages, the resist pattern was transferred to the heterostructure by wet chemical etching with an etch depth of 160 nm. The samples were measured after a brief illumination. The measurements were made in a pumped ^{3}He system with temperatures down to 300 mK using a standard lock-in technique operating at 13 Hz with a constant excitation current of $I_{ex} \leqslant 10$ nA.

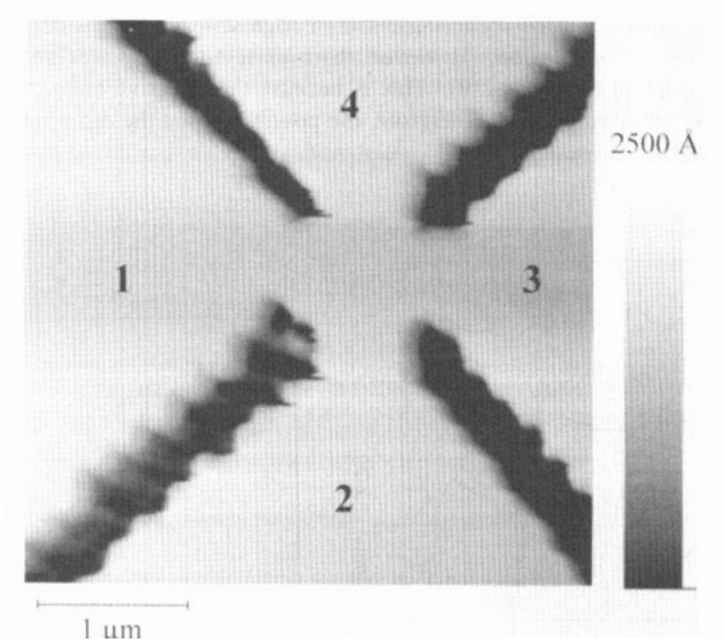

Figure 1. An atomic force-microscopy image of a $W_{\text{lith}} = 0.8$ μ m sample.

3. Experimental results

Figure 2 shows the bend resistance $R_B = R_{12,43}$, where $R_{kl,mn}$ is defined as a four-terminal resistance with the current flowing from lead k to lead l and the voltage difference being measured between leads m and n, for two different conduction widths. The dip at $B = 0$ in the $W_{\text{lith}} = 0.8$ μm sample is due to ballistic transport of electrons. According to the Büttiker–Landauer formula [16], R_B in a symmetric cross at $B = 0$ is described as

$$R_B = -\left(h/e^2\right)\ (T_F - T_S)\,/4T_S\,(T_F + T_S) \tag{2}$$

where T_F is the transmission probability into the front probe, while T_S is the transmission probability into each of the side probes. When the electron transport is ballistic, i.e. $T_F > T_S$, R_B becomes negative [15]. The negative R_B vanishes in the presence of a strong magnetic field, as the electrons are deflected into the side probe favoured by the Lorentz force [17]. In a cross junction of straight wires, the classical transmission probabilities at $B = 0$ can be calculated analytically. Using $T_F/N = 0.414$ and $T_S/N = 0.293$, one obtains $R_B = -3.77$ kΩ N^{-1}, where $N = k_F W/\pi$ is the number of occupied subbands in the wire of width W and k_F is the Fermi wavevector. Within this assumption, N is roughly estimated to be ~ 12 from the amplitude of the negative R_B near $B = 0$, giving rise to $W \sim 0.4$ μm. When B is increased, R_B shows minima corresponding to IQHE and FQHE. As indicated by an arrow in figure 2(a), a pronounced dip is observed around $\nu = 1/2$ in the narrow sample. The dip was observed in both of two $W_{\rm lith} = 0.8$ μm samples we measured. On the other hand, the large ($W_{\rm lith} = 4.0$ μm) sample exhibits only a shallow dip, which is always observed in the resistance of a 2DEG [18]. This clearly demonstrates that the sharp dip originates from the confinement effect in the junction. As the transport phenomena due to the composite fermions occur at high magnetic fields, one may suspect that the dip near $\nu = 1/2$ is due to complicated scattering between edge states, which is determined by the potential near the cross junction. However, there exist no edge states near $\nu = 1/2$ as reported by Wang and Goldman [19]. This is because the effective magnetic field for the composite fermions is very small. Therefore, the possibility that the dip originates from the inter-edge scattering is excluded.

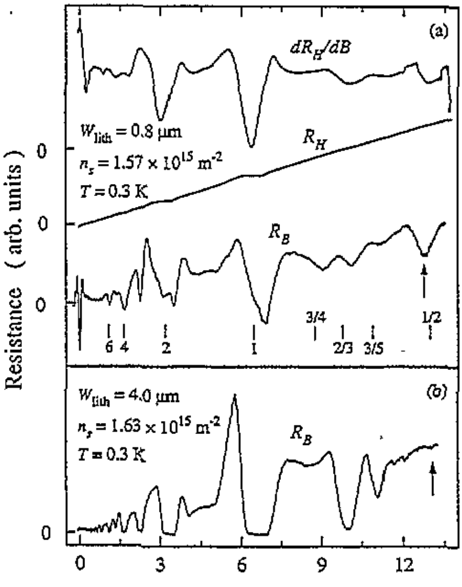

Figure 2. The bend resistance R_B, Hall resistance R_H and its derivative dR_H/dB in (*a*) narrow and (*b*) wide cross junction. The arrows indicate the dips near $\nu = 1/2$. The vertical bars indicate the positions of the fractions calculated from the electron density.

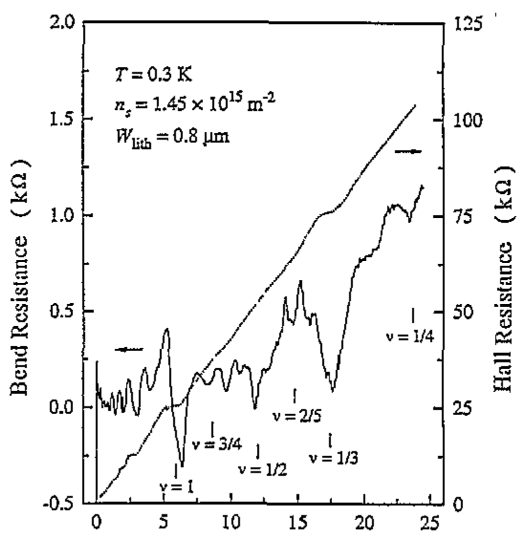

Figure 3. The bend resistance and Hall resistance in a narrow cross junction. The vertical bars indicate the position of the fractions.

Figure 2(a) also shows $R_H = R_{13,24}$ and its derivative with respect to B in the narrow junction. At weak magnetic field of $B \sim 250$ mT, R_H exhibits the 'last plateau' [14]. For lower n_s (smaller W), the quenching is also observed in R_H though it disappears for the n_s shown in figure 2(a). In higher magnetic fields, R_H is quantized at approximately $(h/e^2)/\nu$ for $\nu = 1$ and 2. One can also resolve the plateaus associated with higher integer ν. Since the subbands are defined by both magnetic and electrostatic confinement, the Landau plot is found to show a deviation from the linear behaviour [20]. The effective conduction width estimated from the deviation is in reasonable agreement with the above mentioned estimate of 0.4 μm. Beyond $B = 7$ T, there are features associated with the FQHE, which correspond to the minima in R_B. Some dips in R_B (for instance, for $\nu = 1$) are shifted to higher B relative to the plateaus in R_H. This is ascribed to inter-boundary scattering between edge states in the cross region. When the topmost edge state is nearly depleted, the scattering between the states at opposite boundaries of the channel becomes significant. This shifts the dips in R_B to the higher-magnetic-field side of the quantum Hall states. Around $\nu = 1/2$, the slope of R_H becomes smaller than the linear dependence in a 2DEG. The Hall resistance in a conventional 2DEG is usually featureless around $\nu = 1/2$. When reducing the constriction width by decreasing the amount of the illuminations, the anomalies are, in principle, expected to become stronger. However, additional resistance fluctuations, plausibly due to the inter-boundary scattering, grow rapidly in narrow samples. The mean free path is also suppressed drastically for lower n_s, giving rise to conductance fluctuations. Therefore, we show only the results for large n_s, for which we can unambiguously identify the structures originating from the IQHE and the FQHE.

In figure 3, we show R_B and R_H up to $B = 25$ T. One finds minima in R_B near $\nu = 1/3$ and 2/5 and a well-developed plateau in R_H near $\nu = 1/3$ due to the FQHE. In addition, we find a dip near $\nu = 1/4$, which we ascribe to ballistic transport of CFs with $s = 2$. The dips associated with the $\nu = 1$ and 1/3 quantum Hall states are shifted to higher B, while those near $\nu = 1/2$ and 1/4 are shifted to lower B, indicating their different origin. The dip at $B \sim 8.5$ T may be due to the metallic state at $\nu = 3/4$, which is the particle–hole conjugation of the $\nu = 1/4$ state. The observation of dips at $\nu = 1/2$, 1/4 and 3/4 provides strong evidence of the ballistic transport of the CFs. Notice that the amplitude of the dips near $\nu = 1/4$ and 3/4 is smaller than that near $\nu = 1/2$, indicating stronger scattering for the CFs with $s = 2$ from the Chern–Simons gauge field. The dips become broader in B as s is increased.

The similarity of the structures at $B = 0$ and $\nu = 1/2$ indicates that those at $\nu = 1/2$ can be explained in terms of ballistic transport of CFs. However, we have observed several different behaviours of the CFs. The negative R_B at $B = 0$ [17] is known to have very weak temperature dependence and has been observed at room temperature in an $In_xGa_{1-x}As/Al_yGa_{1-y}As$ system [21]. The broad dip at $\nu = 1/2$ which is ascribed to a Fermi liquid state of the CF becomes independent of temperature below 1 K and persists up to 10 K [18]. One can anticipate a weak temperature dependence for the dip at $\nu = 1/2$ if it originates from the classical ballistic transmission of the CFs. Figure 4 shows R_B for three different temperatures. The negative R_B at $B = 0$ is almost independent of T in the range shown here and vanishes around $T \sim 40$ K. On the other hand, a strong temperature dependence is found for the dip at $\nu = 1/2$. As shown in the inset of figure 4, the dip almost disappears above $T \sim 3$ K, which is one order of magnitude lower than that for the electrons. The rather strong temperature dependence has been reported in the previous ballistic transport experiments of the CFs in microstructures [9, 10]. Because of the large effective mass of CFs [6, 9], the Fermi energy of the CFs is smaller than that of the electrons. The thermal energy k_BT becomes appreciable at lower T for the CFs. Morf and d'Ambrumenil [22] estimated the energies relevant to the binding of flux and the formation of a Fermi sea. The energy has been found to scale as $n_s^{1/2}$ and is ~ 4 K in our samples, in agreement with our observations. In principle, one can distinguish the FQHE and the negative R_B since the resistance approaches zero but is always positive as T is lowered for the FQHE, whereas it can be negative for the R_B. However, the presence of a monotonic background resistance prevents us from identifying the origin of the dip by its polarity. The background resistance was found to change as the current excitation was varied or after thermal cycles (as seen in figures 2 and 4). However, the amplitude of the dips was confirmed to remain unchanged.

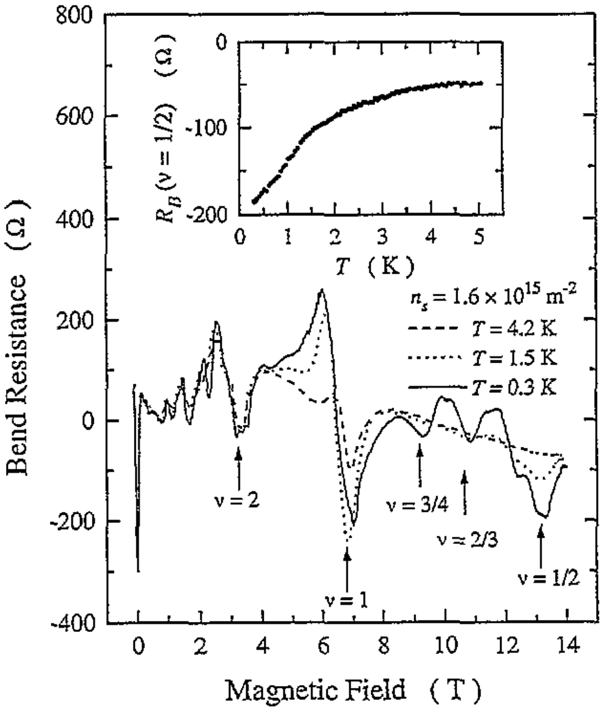

Figure 4. The bend resistance in a $W_{lith} = 0.8$ μm sample for three different temperatures. The dip at $B = 0$ is nearly unchanged. The arrows indicate the dips near $\nu = 2$, 1, 3/4, 2/3 and 1/2. Inset: the bend resistance at $\nu = 1/2$ as a function of temperature.

4. Discussion

The effective mass of the CFs is several times larger than that of the electrons [5–7]. Consequently, the mean free path of the CFs l_0^{CF} is considerably smaller than that of the electrons l_0^{el}. A comparison of the mean free paths, l_0^{el} at $B = 0$ and l_0^{CF} at $\nu = 1/2$, is made in figure 5. We have taken into account the factor of $\sqrt{2}$ due to the spin polarization in l_0^{CF} [5]. As n_s is increased, l_0^{el} increases linearly. The CF mean free path is less than or comparable to the effective constriction width, and so the ballistic transmission probability for the CFs should be significantly reduced. However, the amplitude of the dip at $\nu = 1/2$ is suppressed only by a factor of 2–3 compared to that at $B = 0$. The CFs have been reported to experience short-range scattering [6]. Therefore, the ratio of the ballistic mean free paths of electrons and CFs is smaller than l_0^{el}/l_0^{CF}, which will partly account for the discrepancy. Our results may indicate that the effective constriction width for the CFs is narrower than that for the electrons [10]. We will return to this point later. As reported by Du *et al* [7], the CF mass diverges as $\nu \to 1/2$, implying that the scattering time becomes infinitely long to give a finite resistance. Therefore, the ordinary mean free path may no longer be meaningful for $\nu = 1/2$.

As the Fermi wavevector for the CFs is $k_F^{CF} = \sqrt{4\pi n_s}$, the magnetic field scale will be expanded by a factor of $\sqrt{2}$ if the structures near $\nu = 1/2$ are due to the CFs [5]. However, the dip near $\nu = 1/2$ is much wider in B than that near $B = 0$. The confinement of carriers in the samples is achieved by lateral depletion. The potential near the boundary of the channel varies gradually. Therefore, the carrier density across the channel is not uniform (at least from a classical point of view). This may result in an average over B for the magnetoresistance of the CFs since $\Delta B = 0$ is realized at lower B for the CFs away from the centre of the channel [23]. The non-uniform effective magnetic field explains the shift of the dips to lower magnetic fields. It is also implied that the channel width for the electrons and the CFs are not necessarily the same. The transverse confinement may be enhanced by the finite effective magnetic field near the channel boundary. The effect may be important when one intends to carry out a conductance quantization experiment of the CFs in quasi-one-dimensional systems.

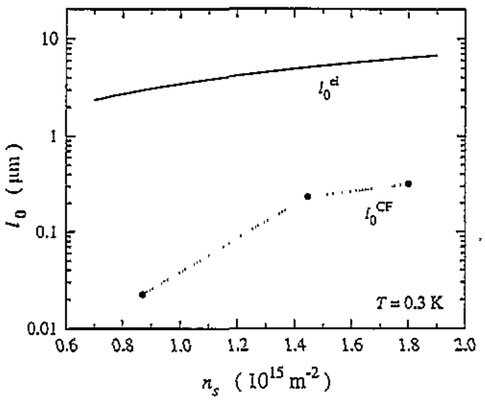

Figure 5. The transport mean free paths l_0 of the electrons and of the composite fermions in unpatterned heterostructure as a function of carrier density n_s.

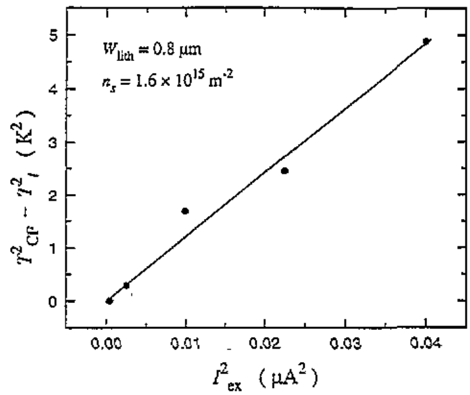

Figure 6. The composite fermion temperature T_{CF} deduced from the amplitude of the negative resistance near $\nu = 1/2$ as a function of the excitation current I_{ex}. The lattice temperature T_l is 0.3 K. The line is a guide to the eye.

As I_{ex} exceeds 20 nA, the amplitude of the dip at $\nu = 1/2$ is suppressed due to the heating effect. The strong temperature dependence of the dip allows us to estimate the CF temperature T_{CF} from the amplitude. The energy balance leads to a relation

$$\tau_e \sigma E^2 = \int_{T_l}^{T_{CF}} c(T)\,\mathrm{d}T \propto T_{CF}^2 - T_l^2 \qquad (3)$$

where τ_e is the energy loss time, E the electric field, σ the conductivity, and $T_l (= 0.3$ K) the lattice temperature. At low temperatures, the specific heat at $\nu = 1/2$ is $c(T) = \pi m k_B^2 T/6$ [5]. As shown in figure 6, equation (3) gives good qualitative agreement with the experimental result.

In conclusion, we have investigated the bend resistance in narrow crossed-wire junctions in the fractional quantum Hall regime. The negative resistance dips near $\nu = 1/2$ and 1/4 are ascribed to the ballistic transport effect of the CFs. The large amplitude of the dips and the width in magnetic field suggest that the effective channel width for the CFs is narrower than that for electrons at zero magnetic field. The channel width narrowing is anticipated to occur due to a nonuniform effective magnetic field for the CFs in the transverse direction.

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References

[1] Tsui D C, Störmer H L and Gossard A C 1982 *Phys. Rev. Lett.* **48** 1559
[2] Laughlin R B 1983 *Phys. Rev. Lett.* **50** 1395
[3] Haldane F D M 1983 *Phys. Rev. Lett.* **51** 605
Halperin B I 1984 *Phys. Rev. Lett.* **52** 1583
[4] Jain J K 1989 *Phys. Rev. Lett.* **63** 199
[5] Halperin B I, Lee P A and Read N 1993 *Phys. Rev.* B **47** 7312
[6] Du R R, Störmer H L, Tsui D C, Pfeiffer L N and West K W 1993 *Phys. Rev. Lett.* **70** 2944
Leadley D R, Nicholas R J, Foxon C T and Harris J J 1994 *Phys. Rev. Lett.* **72** 1906
[7] Du R R, Störmer H L, Tsui D C, Pfeiffer L N and West K W 1994 *Phys. Rev. Lett.* **73** 3274
[8] Willett R L, Ruel R R, West K W and Pfeiffer L N 1993 *Phys. Rev. Lett.* **71** 3846
[9] Kang W, Störmer H L, Pfeiffer L N, Baldwin K W and West K W 1993 *Phys. Rev. Lett.* **71** 3850
[10] Goldman V J, Su B and Jain J K 1994 *Phys. Rev. Lett.* **72** 2065
[11] Suen Y W, Engel L W, Santos M B, Shayegan M and Tsui D C 1992 *Phys. Rev. Lett.* **68** 1379
Eisenstein J P, Boebinger G S, Pfeiffer L N, West K W and Song He 1992 *Phys. Rev. Lett.* **68** 1383
[12] Timp G, Behringer R, Cunningham J E and Howard R E 1989 *Phys. Rev. Lett.* **63** 2268
[13] Chui S T 1986 *Phys. Rev. Lett.* **56** 3011
Chui S T 1987 *Phys. Rev.* B **36** 2806
[14] Roukes M L, Scherer A, Allen S J, Craighead H G, Ruthen R M, Beeby E D and Harbison J P 1987 *Phys. Rev. Lett.* **59** 2395
[15] Takagaki Y, Gamo K, Namba S, Ishida S, Takaoka S, Murase K, Ishibashi K and Aoyagi Y 1988 *Solid State Commun.* **68** 1051
[16] Büttiker M 1986 *Phys. Rev. Lett.* **57** 1761
[17] Takagaki Y, Gamo K, Namba S, Takaoka S, Murase K and Ishida S 1989 *Solid State Commun.* **71** 809
[18] Jiang H W, Störmer H L, Tsui D C, Pfeiffer L N and West K W 1989 *Phys. Rev.* B **40** 12013
[19] Wang K W and Goldman V J 1991 *Phys. Rev. Lett.* **67** 749
[20] Berggren K F, Thornton T J, Newson P I and Pepper M 1986 *Phys. Rev. Lett.* **57** 1769
[21] Hirayama Y and Tarucha S 1993 *Appl. Phys. Lett.* **63** 2366
[22] Morf R and d'Ambrumenil N 1995 *Phys. Rev. Lett.* **74** 5116
[23] Brey L 1994 *Phys. Rev.* B **50** 11861
Chklovskii D B 1995 *Phys. Rev.* B **51** 9895